

THE INTERACTION OF COSMIC RAYS WITH DIFFUSE CLOUDS

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ABSTRACT

We study the change in cosmic-ray pressure, the change in cosmic-ray density, and the level of cosmic-ray induced heating via Alfvén-wave damping when cosmic rays move from a hot ionized plasma to a cool cloud embedded in that plasma. The general analysis method outlined here can apply to diffuse clouds in either the ionized interstellar medium or in galactic winds. We introduce a general-purpose model of cosmic-ray diffusion building upon the hydrodynamic approximation for cosmic rays (from McKenzie & Völk and Breitschwerdt and collaborators). Our improved method self-consistently derives the cosmic-ray flux and diffusivity under the assumption that the streaming instability is the dominant mechanism for setting the cosmic-ray flux and diffusion. We find that, as expected, cosmic rays do not couple to gas within cool clouds (cosmic rays exert no forces inside of cool clouds), that the cosmic-ray density does not increase within clouds (it may slightly decrease in general, and decrease by an order of magnitude in some cases), and that cosmic-ray heating (via Alfvén-wave damping and not collisional effects as for ~ 10 MeV cosmic rays) is only important under the conditions of relatively strong ($10 \mu\text{G}$) magnetic fields or high cosmic-ray pressure ($\sim 10^{-11}$ ergs cm^{-3}).

1. INTRODUCTION

The possibility of a rise or fall in cosmic-ray density within diffuse clouds and/or molecular cores has been studied over many years by researchers interested either in gamma-ray emission from cool clouds (e.g., Skilling 1971; Skilling & Strong 1976; Cesarsky & Voelk 1977; Strong & Skilling 1977; Cesarsky & Völk 1978; Morfill 1982a,b; Voelk 1983) or in ionization by cosmic rays within clouds (e.g., Hartquist et al. 1978; Suchkov et al. 1993; Padoan & Scalo 2005; Padovani et al. 2009; Papadopoulos 2010). Another regime where this interaction may be important is in cool clouds that are readily observed within galactic winds (e.g., Heckman et al. 1990; Martin 2005; Rupke et al. 2005; Steidel et al. 2010). In the context of such winds, cosmic rays are known to help heat and drive outflows of fully ionized gas (Breitschwerdt et al. 1991; Everett et al. 2008; Socrates et al. 2008; Everett et al. 2010); can they help heat and accelerate cool clouds embedded in a hot wind?

We will examine this problem using a set of hydrodynamic equations that treat both convective and diffusive transport of cosmic rays self-consistently; this has rarely been done in studies of cosmic-ray dynamics, and has never been done in the context of cool clouds. Before we start, however, it is important to understand the variety of previous work on this problem. We start, therefore, with a review of models of cosmic-ray interaction with cool clouds in Section 2; we cover models of cosmic-rays in the ISM (Section 2.1) & cosmic rays in galactic winds (Section 2.2), followed by an introduction to our method (Section 2.3). Section 3 defines our equations for cosmic-ray hydrodynamics, and Section 4 applies those models to a simple hot-gas/cool-cloud interface, and examines how robust those results are. Finally, Section 5 summarizes our results and compares & contrasts those results with previous work.

2. REVIEW OF PREVIOUS STUDIES OF COSMIC RAYS IN INHOMOGENEOUS MEDIA

2.1. Cosmic Rays in the Multiphase Interstellar Medium

Galactic cosmic rays are thought to be energized in shocks, most likely in supernova remnants, but perhaps in other shocks as well (see, e.g., Berezhinskii et al. 1990; Schlickeiser 2002; Hillas 2006). Cosmic rays then propagate outward along magnetic field lines, moving through the interstellar medium (ISM). Cosmic rays gyrate around these field lines, but also change direction, scattering off of already-existing magnetic-field inhomogeneities (such as magnetic mirrors, e.g., Cesarsky & Völk 1978; Chandran 2000a), ambient turbulence (e.g., Jokipii 1966; Chandran 2000b; Yan & Lazarian 2004), and Alfvén waves generated resonantly by the cosmic rays themselves (the streaming instability; see Lerche 1966, 1967; Wentzel 1968; Kulsrud & Pearce 1969; Tademaru 1969; Wentzel 1969; Skilling 1970, 1971, 1975; Kulsrud 2005). In this work, and in the review in this section, we will concentrate on cosmic-ray scattering due to turbulence generated by cosmic rays themselves. Not only is the streaming instability a leading candidate for the propagation of $\sim 1 - 100$ GeV cosmic rays, but focusing on the streaming instability helps to demarcate a relatively well-defined problem where we can self-consistently define a cosmic-ray flux and diffusivity. Also, the streaming instability is a key component of any study of the interaction of cosmic rays with magnetic fields in the hot ISM (see also Cesarsky & Völk 1978, who showed that no other turbulent-wave energy source would lead to significant scattering within clouds).

First, we give a short overview of the streaming instability (see also Lerche 1966; Wentzel 1968; Kulsrud & Pearce 1969; Tademaru 1969; Wentzel 1969; Skilling 1971; Zweibel 2003; Kulsrud 2005; Everett et al. 2008); understanding this instability

will be essential in understanding previous results. Simply put, the streaming instability feeds off of the drifting motion of the cosmic-ray population: for instance, the cosmic rays may drift as they stream down their own density gradient. The streaming instability is launched if they move (in bulk) more quickly than the local Alfvén speed; in this case, their cyclotron motion around magnetic-field lines excites Alfvén waves with wavelength of order the cosmic-ray gyroradius, as they scatter off of Alfvén waves. If those waves are not damped much more quickly than they are excited, the Alfvén waves will then scatter the cosmic rays, limiting their bulk motion to approximately the local Alfvén speed. In applying this mechanism to the ionized ISM, the amplitude of the Alfvén waves can be quite small ($\delta B \sim 10^{-3} B_{\text{ambient}}$); with an Alfvén wave of this strength, the cosmic-ray mean-free path is $\lambda_{\text{mfp}} \sim 0.1$ pc (see Kulsrud 2005). The excitation time-scale for these waves is of order 10^4 years in the ISM, so as long as the damping rates are not faster than this, this instability will be important. Of course, as we show in this paper, as the balance between wave damping and wave growth changes, the cosmic-ray mean-free path varies strongly from $\lesssim 0.01$ pc to larger than the size of diffuse clouds; this will be shown most explicitly by the variation in κ_{cr} , explored in Section 4.4.1.

The study of the streaming instability and cosmic-ray propagation from ionized plasma into diffuse clouds started with Kulsrud & Pearce (1969), who pointed out the importance of ion-neutral damping to quickly destroy Alfvén waves generated by the streaming instability (see also Kulsrud & Cesarsky 1971). They inferred from this that one could ignore the interstellar clouds for the purposes of cosmic-ray diffusion: in their work, cosmic rays move at approximately the Alfvén speed in the ionized regions of the ISM only, and free-stream at $v \sim c$ in molecular clouds. A similar result was obtained by Skilling (1971), who named such regions “free zones”: these are regions where cosmic rays are not locked to the Alfvén speed because of wave damping.

After these initial studies, interest in gamma-ray emission from clouds brought this problem back to prominence, with Skilling & Strong (1976) and Cesarsky & Voelk (1977) further investigating how cosmic rays interact with clouds (see also Cesarsky et al. 1977; Strong & Skilling 1977; Cesarsky & Völk 1978; Hartman et al. 1979; Voelk 1983). Skilling & Strong (1976) realized that the destruction of low-energy ($E \lesssim 300$ MeV) cosmic-ray protons within the cloud, due to ionization and nuclear interactions, causes a decrease in the cosmic-ray density at these low energies. This means that the cosmic-ray flux exiting a cloud is lower than the input cosmic-ray flux, and this flux difference excites the streaming instability: cosmic-ray generated Alfvén waves grow in a cosmic-ray density gradient. The streaming instability then generates Alfvén waves, which act as scattering centers for these low-energy cosmic rays, slowing down these cosmic rays from their free-streaming velocity to approximately the Alfvén speed, v_A . This decreases the rate at which cosmic rays can ‘flow’ into the cloud, or put another way, it slows the rate at which cosmic rays are replenished within the cloud; Skilling & Strong (1976) referred to this process as therefore “excluding” cosmic rays from the cloud. Cesarsky & Voelk (1977) studied this mechanism (as well as a variety of others), finding that this exclusion would occur only for cosmic rays with $E \lesssim 50$ MeV if the magnetic field strength in the cloud increased to $50 \mu\text{G}$ as opposed to the constant $3 \mu\text{G}$ assumed in Skilling & Strong (1976); the compression of the magnetic field raises the cosmic-ray flux into the cloud, and as a result, only the lowest-energy cosmic rays (again, $E \lesssim 50$ MeV) are not replenished quickly enough in the cloud (Cesarsky & Voelk 1977; Voelk 1983). (We note that very recent observations of magnetic field strength in cool phases of the ISM seem to show that the magnetic field does not increase significantly in strength for $n \lesssim 300 \text{ cm}^{-3}$; see Crutcher et al. 2010.) It is important to note that these previous analyses seem to rely upon a uniform distribution of cosmic rays far away from the cloud, so that the streaming instability only operates in the ionized region on the periphery of the clouds (also, the importance of the interaction of multiple clouds is briefly addressed in Morfill 1982b).

We mention briefly, for completeness, that cosmic-ray electrons could also be impacted in these processes. Cesarsky & Völk (1978) pointed out that scattering of cosmic-ray electrons off of the Alfvén waves generated by cosmic-ray protons with the same gyro-radii would help slow cosmic-ray electrons within clouds. Morfill (1982a) and Morfill (1982b) expanded on this idea by considering how cosmic-ray electrons might be accelerated within the clouds and how secondaries would be generated within the cloud, helping to increase the γ -ray brehmßtrahlung emission.

But again, the above mechanisms operate on cosmic rays at relatively low energies (of order 50 MeV or lower), and so would not exclude ~ 1 GeV cosmic rays from clouds. This is important for this study, as ~ 1 GeV cosmic rays are the dominant source of cosmic-ray pressure, which is our prime interest here. So, while this “exclusion” effect may occur in the ISM, and this mechanism would be important for understanding the ionization of cool clouds by cosmic rays (Hartquist et al. 1978; Suchkov et al. 1993; Padovani et al. 2009; Papadopoulos 2010), it will not strongly impact cosmic-ray pressure or GeV gamma-ray emission.

After the above papers were published, it seems that cosmic rays were then left to stream on their own for a few years. However, in the past decade, as cosmic-ray ionization rates have become better constrained (e.g., McCall et al. 2003; Padovani et al. 2009; Indriolo et al. 2010a,b), and as observations have started to focus on interactions of cosmic rays with molecular clouds (e.g., Aharonian 2001; Gabici et al. 2007; Protheroe et al. 2008; Torres et al. 2005, 2008; Li & Chen 2010; Gabici et al. 2010; Fatuzzo et al. 2010; Casanova et al. 2011; Uchiyama 2011; Ohira et al. 2011; Ellison & Bykov 2011), interest in understanding the interaction of clouds and cosmic rays has increased. However, almost all of the above works consider parameterized cosmic-ray diffusion, in sharp contrast to the models of the late-70s and early-80s, which worked to understand cosmic-ray diffusion by studying the streaming instability.

An exception to this trend is the work of Padoan & Scalo (2005). These authors calculated that, as cosmic rays move from the hot, ionized medium into diffuse clouds, the cosmic-ray density strongly increases (with $n_{\text{cr}} \propto n_i^{1/2}$ up to approximately $n_i \sim 100 \text{ cm}^{-3}$); they also derived that n_{cr} is approximately constant in molecular clouds of higher densities. Their work is a response to the observations of McCall et al. (2003), who infer an increase in the cosmic-ray-induced ionization rate with clouds to explain observations of H_3^+ . The theoretical result of Padoan & Scalo (2005) seems to contradict, however, the above-mentioned, earlier studies which showed an effect only for low-energy cosmic rays, and predicted only a *decrease* in cosmic-ray density at those energies. How do we understand this difference? How do we predict the pressure within clouds, or the gamma-ray flux from diffuse clouds and molecular clouds?

2.2. Cosmic Rays in Multiphase Galactic Winds

Galactic winds also have multiple gas phases, just like the ISM that launches those winds. For instance, Na I and Mg II absorbers have been detected in the halos of luminous and ultra-luminous infrared galaxies; these neutral and near-neutral components are apparently outflowing from their host galaxies with velocities of hundreds of kilometers per second (Heckman et al. 1990, 2000; Shapley et al. 2003; Martin 2005; Rupke et al. 2005; Steidel et al. 2010). The clouds (observed to have small line-of-sight covering fraction; see Chen et al. 2010) are important for a number of reasons. First, the velocities of cool clouds are much easier to measure than the velocity of the hot-gas component in winds, and so such clouds can be used to constrain the velocities of the hot wind (if we understand how the clouds interact with the hot wind). Second, it is hard to understand how such clouds survive (or perhaps re-form?) as coherent entities in the wind for timescales approaching tens of millions of years (Marcolini et al. 2005; Cooper et al. 2008, 2009; Murray et al. 2010). Interestingly, to our knowledge, there has been no consideration of how cosmic rays (which could help drive a galactic wind; see Ipavich 1975; Breitschwerdt et al. 1991; Everett et al. 2008; Socrates et al. 2008) would interact with clouds in galactic winds.

We point out that previous work has considered other driving mechanisms for such cool clouds: in particular, much research on this problem has focused on the radiative driving of clouds via radiative acceleration on dust in the clouds (e.g., Nath & Silk 2009; Murray et al. 2010), or perhaps by ram pressure from the surrounding hot-phase component of the wind (Murray et al. 2005, 2010). However, since cosmic rays are already important in winds, could they also help accelerate clouds? And, since cosmic rays can penetrate throughout each cloud, could they help to drive the entire mass of a cloud (instead of only pushing on the boundaries) without introducing instabilities on the cloud boundary that could destroy it?

2.3. An Overview of Our Approach To Both Problems

For this paper, we will concentrate on the general interaction of ~ 1 GeV cosmic rays with clouds; we will consider parameters applicable both to diffuse clouds embedded in the $\sim 10^6$ K, ionized ISM as well as cool clouds in galactic winds. We do not explicitly treat the interaction with molecular clouds; however, as we will see, cosmic rays already start to freely stream (without scattering) in diffuse clouds and their behavior will not change qualitatively in molecular clouds, with the exception that cosmic-ray losses will become more extreme at high density. In this sense, as long as our assumptions of, for instance, magnetic-field geometry hold, our work can be extended to molecular clouds as well.

We find that it is relatively straightforward to investigate these problems (under some important restrictions) using the equations of cosmic-ray hydrodynamics (McKenzie & Voelk 1981, 1982; McKenzie & Webb 1984; Breitschwerdt et al. 1991; Zirakashvili et al. 1996); we have used similar equations in previous papers (Everett et al. 2008, 2010). These equations describe the interaction of cosmic rays with an ambient magnetized medium through the action of the streaming instability; as long as the cosmic-ray mean-free path is larger than the scales of interest in the problem, we can use these equations to model the cosmic-ray pressure, Alfvén wave pressure, and the diffusivity of cosmic rays. These equations also assume that the cosmic-rays are scattered strongly enough that they are nearly isotropic in pitch angle (the angle between the cosmic-ray velocity vector and the magnetic field); for the strong scattering in the ionized phase of the ISM, and for cosmic rays of energy near 1 GeV, both of these requirements are easily satisfied.

For the case of clouds in the ISM, we will investigate the cosmic-ray pressure and density in the cloud as a function of the boundary conditions for those diffuse clouds. We will ask whether there is an appreciable cosmic-ray pressure gradient in such clouds for cosmic rays with $E \sim 1$ GeV, and whether the heating due to damped Alfvén waves on the periphery of the cloud is significant.

We will ask the same questions for cosmic rays interacting with clouds within a galactic wind: what is the cosmic-ray pressure within clouds? If the cosmic-ray density or pressure actually increases within clouds as postulated by Padoan & Scalo (2005), can cosmic rays help destroy the cloud? Will the cosmic rays spark significant heating within the cloud or at least at the cloud boundary? As mentioned above, if these clouds are to be used to understand mass outflow from galaxies (especially as it relates to a surrounding, hot outflowing wind component), understanding the acceleration mechanism of these clouds, and their interaction with all of the components of the hot wind around them, are both very important.

Before explaining the setup of our models, we stress that, of course, the idea of a “cloud” itself is an idealization, as is a particular structure for the cloud and its interface with the ambient medium. Given the uncertainties here, we do not intend to present a study of the interactions of cosmic rays with any particular wind- or ISM-cloud structure; our goal is to understand how cosmic rays interact with generic cool clouds under a range of simplified assumptions, and to therefore understand generically what role cosmic rays can play in this multiphase medium, and to point the way to observational tests of the role of cosmic rays in a variety of multiphase media.

3. COSMIC-RAY HYDRODYNAMICS

When writing the equations of cosmic-ray hydrodynamics, researchers usually assume that there is zero cosmic-ray diffusivity with respect to Alfvén waves (e.g., Breitschwerdt et al. 1991; Everett et al. 2008). In the case of cool clouds, however, it is known that ion-neutral friction in the cool clouds will damp the cosmic-ray generated Alfvén waves so strongly (see, e.g., Zweibel & Shull 1982; Balsara 1996) that it is possible for cosmic rays to be very loosely coupled to the Alfvén waves, and diffuse relative to the waves (e.g., Kulsrud & Cesarsky 1971). Therefore, for this work, we utilize the cosmic-ray hydrodynamic equations (as derived and presented in McKenzie & Voelk 1981, 1982; McKenzie & Webb 1984; Breitschwerdt et al. 1991; Zirakashvili et al. 1996; Ptuskin et al. 1997), retaining all of the usual terms in the above papers, and also including a diffusion term, where the diffusion is set self-consistently by the Alfvén-wave pressure in the medium (see also Ptuskin et al. 1997, who included a diffusion term in their calculation of wind dynamics).

In all of our equations, we use Gaussian-cgs units.

3.1. Cosmic-Ray Hydrodynamics Equations With Diffusion

We start with the cosmic-ray transport equation (Eq. A8 from Breitschwerdt et al. 1991):

$$\frac{\partial}{\partial t} \left(\frac{P_{\text{cr}}}{\gamma_{\text{cr}} - 1} \right) + \nabla \cdot \left(\frac{\gamma_{\text{cr}}}{\gamma_{\text{cr}} - 1} (\mathbf{u} + \mathbf{v}_A) P_{\text{cr}} - \frac{\kappa_{\text{cr}}}{\gamma_{\text{cr}} - 1} \nabla P_{\text{cr}} \right) = (\mathbf{u} + \mathbf{v}_A) \nabla P_{\text{cr}} + Q, \quad (1)$$

where P_{cr} is the cosmic-ray pressure, γ_{cr} is the adiabatic index for cosmic rays, \mathbf{u} is the velocity of the thermal plasma, \mathbf{v}_A is the Alfvén speed, κ_{cr} is the cosmic-ray diffusivity, and Q represents energy input/loss for the cosmic rays. Assuming a time-steady source of cosmic-rays, and evaluating the equation for a quasi-1D system, Equation 1 becomes:

$$\kappa_{\text{cr}} \frac{d^2 P_{\text{cr}}}{dz^2} = \left(\frac{\kappa_{\text{cr}}}{P_w} \frac{dP_w}{dz} + (u + v_A) \right) \frac{dP_{\text{cr}}}{dz} - \frac{\gamma_{\text{cr}} P_{\text{cr}}}{\rho} \left(\left(u + \frac{v_A}{2} \right) - q \right) \frac{d\rho}{dz} - (\gamma_{\text{cr}} - 1) Q, \quad (2)$$

where ρ represents the plasma mass density, and

$$P_w = \frac{\langle (\delta B)^2 \rangle}{8\pi} \quad (3)$$

is the Alfvén-wave pressure. We can simplify this further by setting $u \equiv 0$ and $q \equiv 0$ (no large-scale flows and no plasma mass sources):

$$\kappa_{\text{cr}} \frac{d^2 P_{\text{cr}}}{dz^2} = \left(\frac{\kappa_{\text{cr}}}{P_w} \frac{dP_w}{dz} + v_A \right) \frac{dP_{\text{cr}}}{dz} - \left(\frac{\gamma_{\text{cr}} P_{\text{cr}} v_A}{\rho} \right) \frac{d\rho}{dz} - (\gamma_{\text{cr}} - 1) Q. \quad (4)$$

To derive the above result, we have assumed mass conservation and magnetic-flux conservation. In addition, we have defined the cosmic-ray diffusivity κ_{cr} as

$$\kappa_{\text{cr}} = v \lambda_{\text{mfp}} = \frac{4}{\pi} \frac{c}{3} \frac{r_L}{(\delta B/B)^2}, \quad (5)$$

or, given the definition of the wave pressure, Equation 3, and defining the magnetic pressure as

$$P_B = \frac{B^2}{8\pi}, \quad (6)$$

we can rewrite κ_{cr} as:

$$\kappa_{\text{cr}} = v \lambda_{\text{mfp}} = \frac{4}{\pi} \frac{c}{3} \frac{r_L}{P_w/P_B}. \quad (7)$$

As a check, in the limit of $\kappa \rightarrow 0$, we find

$$\frac{dP_{\text{cr}}}{dz} = \frac{\gamma_{\text{cr}} P_{\text{cr}}}{v_A} \left(\frac{v_A}{2\rho} \frac{d\rho}{dz} \right) + \frac{(\gamma_{\text{cr}} - 1) Q}{v_A}, \quad (8)$$

which, if we further assume $Q = 0$, simplifies to:

$$\frac{dP_{\text{cr}}}{dz} = \frac{\gamma_{\text{cr}} P_{\text{cr}}}{2\rho} \frac{d\rho}{dz}. \quad (9)$$

This equation leads to the useful result, *for cases without diffusion*,

$$P_{\text{cr}} v_A^{\gamma_{\text{cr}}} = \text{constant}, \quad (10)$$

or, taking into account that $\nabla \cdot \mathbf{B} = 0$, we can rewrite this as

$$P_{\text{cr}} \rho^{-\gamma_{\text{cr}}/2} = \text{constant}, \quad (11)$$

which can also be seen from Equation 1.

The Alfvén-wave equation does not require any additional work beyond that presented in Breitschwerdt et al. (1991), except the consideration of the loss term. We start with the full Alfvén-wave equation:

$$\frac{\partial}{\partial t} \left(\frac{(\delta B)^2}{4\pi} \right) + \nabla \cdot \left(\frac{(\delta B)^2}{4\pi} \left[\frac{3}{2} \mathbf{u} + \mathbf{v}_A \right] \right) = \mathbf{u} \nabla P_w - \mathbf{v}_A \nabla P_{\text{cr}} + \mathcal{L}, \quad (12)$$

where \mathcal{L} represents other energy-gain, or, if negative, energy-loss mechanisms for waves. Simplifying this equation under the assumptions of a steady-state, 1D system, we find:

$$\frac{dP_w}{dz} = \frac{3u + v_A}{2(u + v_A)} \frac{P_w}{\rho} \frac{d\rho}{dz} - \frac{v_A}{2(u + v_A)} \frac{dP_{\text{cr}}}{dz} - \frac{3q}{2(u + v_A)} \frac{P_w}{\rho} + \frac{\mathcal{L}}{2(u + v_A)}. \quad (13)$$

In our limit of $u \equiv 0$, $q \equiv 0$, we find:

$$\frac{dP_w}{dz} = \frac{1}{2} \frac{P_w}{\rho} \frac{d\rho}{dz} - \frac{1}{2} \frac{dP_{\text{cr}}}{dz} + \frac{\mathcal{L}}{2v_A}. \quad (14)$$

In the absence of cosmic rays or wave-energy gains or losses, Equation 14 predicts $P_w \propto \rho^{1/2}$, in agreement with McKee & Zweibel (1995).

We then substitute this equation for dP_w/dz into Equation 4. Next, we need to define the wave-pressure loss term, \mathcal{L} .

3.2. Alfvén Wave Loss Mechanisms

This loss term, \mathcal{L} , represents the Alfvén-wave damping mechanisms; we consider both non-linear Landau damping (in the hot, intercloud medium) as derived by Kulsrud (2005):

$$\dot{\mathcal{L}}_{\text{NLLD}} = -2 \cdot P_w \cdot \frac{3}{4} \sqrt{\frac{\pi}{2}} \left(\frac{\delta B}{B} \right)^2 \frac{v_i}{c} \omega_{c,i} \quad (15)$$

(where we have substituted $\omega_{\text{Alfvén}} = kv_A$ with $k \sim \frac{1}{r_g} \sim \frac{c}{3} \frac{1}{\omega_{c,i}}$), and ion-neutral friction (‘case a’ from De Pontieu et al. 2001):

$$\dot{\mathcal{L}}_{\text{INF}} = -2 \cdot P_w \cdot 5 \times 10^{-15} \sqrt{2} v_i \frac{\rho_n}{m_p} \quad (16)$$

where v_i is the ion thermal speed (which we set to $\sqrt{8k_B T / (\pi m_p)}$ throughout this work), δB is the amplitude of the Alfvén-wave perturbation to the large-scale magnetic field (B), and ρ_n is the mass density of neutral atoms. This equation applies for ion-neutral friction between protons and hydrogen atoms; lower cross-sections are necessary for the collision of carbon ions with hydrogen atoms, for instance, in much cooler clouds.

3.3. Energy Loss Mechanisms for Cosmic-Ray Protons

Cosmic-ray protons themselves are subject to energy losses due to pion production, ionization, and Coulomb collisions (e.g., Skilling & Strong 1976); these losses are required in Equation 4, and are represented there with the term \mathcal{Q} . To model these effects as a function of cosmic-ray-proton energy, we utilize a modified version of the the energy loss function given in Schlickeiser (2002) as his Equation (5.3.58):

$$\begin{aligned} -\left(\frac{dT}{dt}\right)_{\text{proton}}(\beta) = & 1.82 \times 10^{-7} \times \left[(n_{\text{HII}} + n_{\text{HI}} + 2n_{\text{H}_2}) \times [5.44(\gamma - 1)^{7.64} \gamma^{0.75} H(1.75 - \gamma) H(\gamma - 1.43) \right. \\ & \left. + 0.72(\gamma - 1)^{0.53} \gamma^{0.75} H(\gamma - 1.75)] \right. \\ & + 1.69 n_e \frac{\beta^2}{\beta^3 + 2.34 \times 10^{-5} (T_e/2 \times 10^6 \text{K})^{3/2}} \\ & \left. \times H[\beta - 7.4 \times 10^{-4} (T_e/2 \times 10^6 \text{K})^{1/2}] \right. \\ & \left. + (n_{\text{HI}} + 2n_{\text{H}_2}) \frac{2\beta^2}{10^{-6} + 2\beta^3} \times [1 + 0.0185(\ln \beta) H[\beta - 0.01]] \right] \text{eV s}^{-1} \quad (17) \end{aligned}$$

In this equation, T is the kinetic energy per proton, $H[\dots]$ is the Heaviside step function, x is the ionization fraction, and $\beta \equiv v/c$. In this expression, the first term describes pion production, the second gives Coulomb losses (normally negligible), and the third term characterizes losses due to ionization. The exact form of this equation in Schlickeiser (2002) made some assumptions valid in the high-energy regime, but which could be problematic at lower energies, so Equation 17 is our modified version of the result from Schlickeiser (2002). In particular, the derivation of the pion-production loss term in Schlickeiser (2002) assumed $\gamma \gg 1$, and is undefined below $\gamma = 1.3$, but we are interested in the regime $\gamma \sim 2$. Equation 17 also more explicitly includes the dependencies on n_{HI} , n_{HII} and n_{H_2} from earlier expressions in Schlickeiser (2002).

Because we are especially interested in the cosmic-ray spectrum near 1 GeV, where the majority of cosmic-ray momentum is, we will start by considering only the energy-loss rate near 1 GeV, but will examine the impact of the variation in energy-loss rate in Section 4.4. This energy loss rate gives the expression for \mathcal{Q} in the equation for the cosmic-ray pressure (Eq. 4).

4. CALCULATING THE COSMIC-RAY PRESSURE IN CLOUDS

4.1. Boundary Conditions and Cloud Setup

To investigate cosmic rays penetrating diffuse clouds using Equations 4 and 14 (with definitions from Eqs. 3, 6, 7, 15, 16, and 17), we set up a simple interface between a highly-ionized plasma and a cool, neutral cloud.

Our first use of this setup will be to check that the code reproduces an analytical result: if there is no cosmic-ray diffusivity (the cosmic rays are perfectly locked to Alfvén waves), the cosmic-ray pressure and Alfvén speed change together to conserve the adiabatic constant $P_{\text{cr}} v_A^{\gamma_{\text{cr}}}$ (Breitschwerdt et al. 1991). To test this, we define the density contrast between the ionized and neutral phases with a simple *tanh* profile (the same profile we will use for all of our calculations in this paper):

$$\rho(z) = \rho_{\text{init}} + (\rho_{\text{final}} - \rho_{\text{init}}) \cdot \frac{1}{2} \left(1 + \tanh \left[\frac{z - z_{\text{cloud edge}}}{\Delta z_{\text{edge}}} \right] \right), \quad (18)$$

with the center of the tanh function at $z_{\text{cloud edge}} = 0.5 \text{ pc}$, a width of $\Delta z_{\text{edge}} = 0.05 \text{ pc}$; the initial and final densities are set to (for this initial test) $\rho_{\text{init}} = 1 m_p \text{ cm}^{-3}$ and $\rho_{\text{final}} = 10 m_p \text{ cm}^{-3}$. We also set the ionization fraction to 1 throughout, and have removed

all loss terms for the cosmic rays. It is not possible to simply set $\kappa_{\text{cr}} \equiv 0$, so we integrate only Equation 8 for dP_{cr}/dz . We then initialize the cosmic-ray pressure to 10^{-12} ergs cm^{-3} , and integrate the equations for cosmic-ray and wave pressure into the cloud. The resultant adiabatic constant $P_{\text{cr}}v_A^{\gamma}$ shows only very small variations; it is constant to 1 part in over 10^4 . This verifies the basic structure of the code.

Next, we define the complete transition from an ionized medium to a cool cloud; this is the basic setup we will explore in this paper. The ionized medium is defined to have $T = 3 \times 10^6$ K, $\rho = 10^{-2} m_p \text{ cm}^{-3}$, $P_{\text{cr}} = 10^{-12}$ ergs cm^{-3} , $B = 3 \mu\text{G}$ (and the same value within the cloud), and x (the ionization fraction) set to unity. As in the previous test, the density is then defined to increase as a simple *tanh* function; the final cloud density is set at $\rho = 100 m_p \text{ cm}^{-3}$. The ionization fraction, x , is set to decrease as:

$$x(z) = x_{\text{final}} + (x_{\text{init}} - x_{\text{final}}) \cdot \frac{1}{2} \left(1 + \tanh \left[\frac{-(z - z_{\text{cloud edge}})}{\Delta z_{\text{edge}}} \right] \right). \quad (19)$$

In this way, the ionization fraction is set to drop from a maximum ($x_{\text{init}} = 1$) to a minimum ($x_{\text{final}} = 10^{-3}$) in a very similar way to the drop in temperature, but lagging slightly behind. Meanwhile, we set the temperature to vary inversely as the density to maintain a constant pressure. The variation in density, ionization fraction and temperature are all shown in Figure 1.

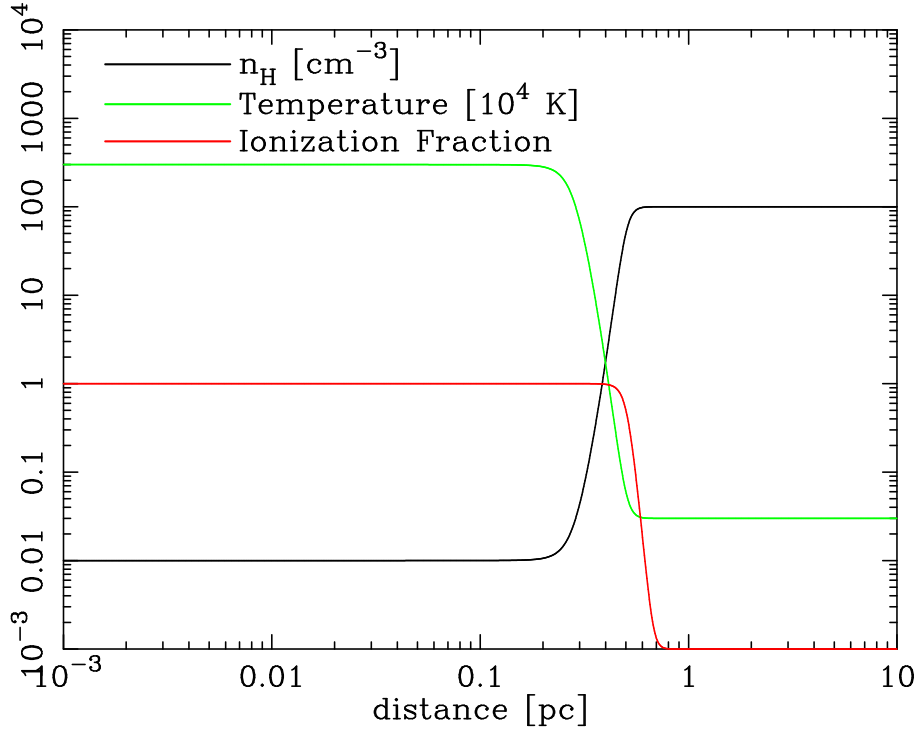


Figure 1. The variations of gas density (in units of cm^{-3}), temperature (in units of 10^4 K), and ionization fraction as cosmic rays transition from the low-density, high-temperature ISM ($x_{\text{init}} = 1$) to a simple cool cloud ($x_{\text{final}} = 10^{-3}$). This transition is not meant to mimic any particular cloud, but is just a fiducial system to examine the change in cosmic-ray pressure.

We set the initial wave pressure such that the initial P_w is in equilibrium with the initial cosmic-ray pressure gradient and the initial wave damping. For our fiducial setup, with a cosmic-ray pressure gradient $dP_{\text{cr}}/dz = -10^{-34}$ ergs cm^{-4} (this value is the approximate cosmic-ray gradient in galactic wind models in Everett et al. 2010), the equilibrium P_w is approximately 2.3×10^{-18} ergs cm^{-3} , yielding an equilibrium cosmic-ray diffusion coefficient of $1.5 \times 10^{27} \text{ cm}^2 \text{ s}^{-1}$ in the ionized medium (or, put in more transparent units, $\kappa_c = 0.005 \text{ kpc}^2/\text{Myr}$). Interestingly, this diffusion coefficient is in the range already considered by some parameterized cosmic-ray diffusion models (e.g., Casanova et al. 2011). However, this value is approximately a factor of 20 to 50 lower than the overall Galactic diffusion coefficient found from fitting diffusion models to cosmic-ray spectra and gamma-ray emission, e.g., Berezhinskii et al. (1990), Ptuskin et al. (2008), and Strong et al. (2010). This difference can be understood as a result of our choice of the input cosmic-ray pressure gradient; a smaller pressure gradient (which could be quite reasonable for P_{cr} differences within the ISM) would lead to a smaller wave pressure, and hence larger diffusivity; in order to explain the difference we see, the cosmic-ray pressure gradient would have to be 2500 times smaller in the ISM, or $dP_{\text{cr}}/dz = -4 \times 10^{-38}$ ergs cm^{-4} . As we report in Section 4.4 (including Table 1), even decreasing the cosmic-ray pressure gradient down to 10^{-39} ergs cm^{-4} (corresponding to $\kappa_{\text{cr}} \sim 5 \times 10^{29} \text{ cm}^2 \text{ s}^{-1}$ or $1.7 \text{ kpc}^2 \text{ Myr}^{-1}$) does not significantly change the results of our fiducial model.

We integrate the diffusion equation for cosmic-ray pressure until the diffusion coefficient (calculated from Eq. 7) exceeds cL_{cloud} (where L_{cloud} is the scale of the cloud, here set to 10 pc), or until P_{cr} drops to zero, which occurs in some extreme cases, as we discuss in Section 4.4. Deeper into the cloud, beyond either of these limits, the cosmic rays are no longer interacting with

waves in the cloud (or in the case where P_{cr} drops to near zero, they are not generating waves in the cloud), and so would enter a purely kinetic regime where we cannot treat them hydrodynamically. Effectively, this happens where-ever ion-neutral damping takes over, damping the cosmic-ray generated waves, and releasing the cosmic rays to free-stream into the cloud.

4.2. Cosmic-Ray Pressure in the Cloud

We show the results for the cosmic-ray pressure, cosmic-ray diffusivity, and Alfvén-wave pressure in our fiducial model in Figure 2. This figure summarizes many of the key findings of the paper: we see a slight (7.5%) decrease in the cosmic-ray pressure, with the coupling of cosmic rays to the gas stopping at a distance of ~ 0.4 pc into the cloud: further inside the cloud, the cosmic rays decouple from the gas (this is indicated by the dotted line in Fig. 2). Beyond that point of decoupling, the cosmic-ray pressure still formally exists within the cloud (the cosmic rays still have a well-defined and positive energy density), but that pressure is not communicated to the gas in the cloud, so there is no cosmic-ray pressure force on the cloud. Interestingly, though, on the outskirts of the cloud, we see that the cosmic-ray diffusivity actually decreases strongly because of a sharp increase in the wave pressure.

How can we understand these effects intuitively? The overall drop in cosmic-ray coupling to near zero is easy to explain as the impact of ion-neutral damping that destroys the Alfvén waves, leading to extremely weak coupling between the cosmic rays and the gas. But why the small decrease in the cosmic-ray pressure beforehand? And why the decrease in the diffusivity on the cloud boundary? For these latter effects, the key is to recognize that the solution is the only self-consistent one that allows the initial cosmic-ray flux to be transported through the cloud without a buildup in cosmic-ray pressure at any point; this will be explained in more detail, below, but can be seen by comparing the advective flux of cosmic rays (the cosmic-ray flux in the ionized medium, where cosmic rays are largely tied to Alfvén waves) to the diffusive flux, as shown in Figure 3.

The incoming advective flux (plotted with the dot-dashed line in Fig. 3; this flux is given by the advection term in Eq. 1, and discussed further after Eq. 21, below) is set by the Alfvén speed and cosmic-ray pressure (the flux is $\gamma_{\text{cr}} v_A P_{\text{cr}}$). As the density increases into the cloud, the Alfvén speed decreases, and so the cosmic-ray flux that can be transported by the waves also decreases, as shown in the downward trend of the dot-dashed line in Figure 3. The only way to avoid a buildup of cosmic-ray pressure on the outskirts of the cloud is to transfer that flux of cosmic rays to a diffusive flux, $\kappa_{\text{cr}} dP_{\text{cr}}/dz$, which depends on the cosmic-ray pressure gradient. At the same time, as an added effect, the increase in gas density leads to an increase in Alfvén-wave pressure (see Eq. 14; see also McKee & Zweibel 1995), and so the diffusion coefficient itself must decrease, requiring a compensating change in the cosmic-ray pressure, P_{cr} . The self-consistent solution is to therefore have a larger cosmic-ray pressure gradient, which results in a dominant diffusive flux into the cloud.

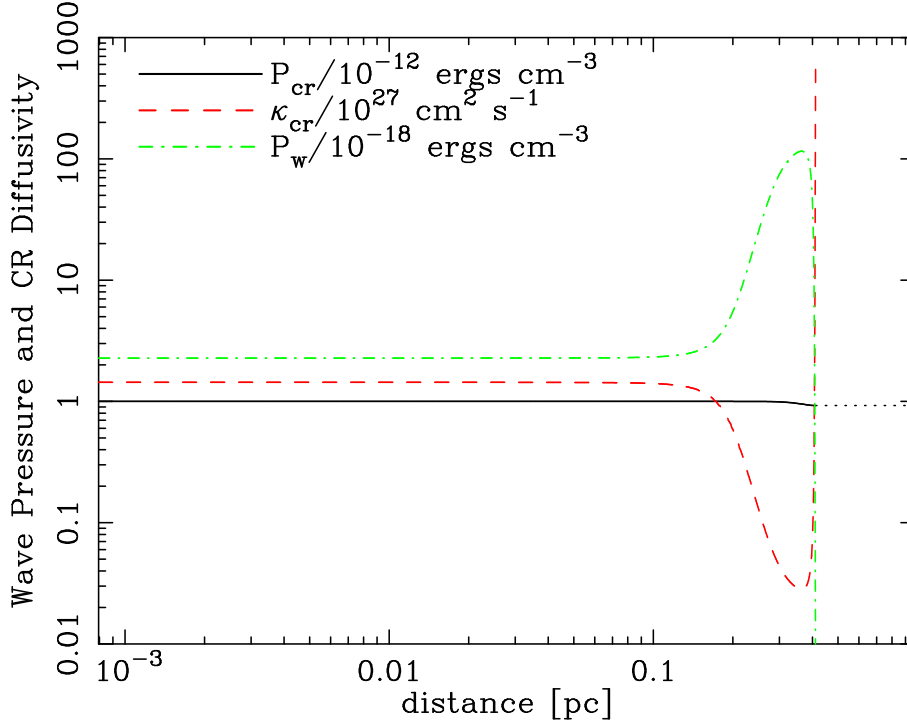


Figure 2. The normalized Alfvén-wave pressure (P_w), cosmic-ray diffusivity (κ_{cr}), and cosmic-ray pressure (P_{cr}) as a function of distance from the external ionized medium (on the left) to the cool cloud (on the right). P_w is initially in equilibrium outside the cloud, then increases on the cloud boundary as the density increases and as the cosmic-ray pressure drops slightly (by of order 5%). Near the peak of P_w , where the growth slows, non-linear Landau damping starts to limit the growth of the waves; the destruction of P_w at $z = 0.4$ pc is completely due to the onset of ion-neutral damping. $\kappa_{\text{cr}} \propto P_w^{-1}$, so closely follows the development of the Alfvén-wave pressure. Beyond $z = 0.4$ pc, κ_{cr} exceeds cL_{cloud} , and so the cosmic rays transition from the hydrodynamic approximation to the kinetic limit; beyond that point, the density of cosmic rays is constant (this is indicated by the dotted line). Our treatment cannot describe cosmic rays in that limit, and so we stop tracking cosmic rays past that point.

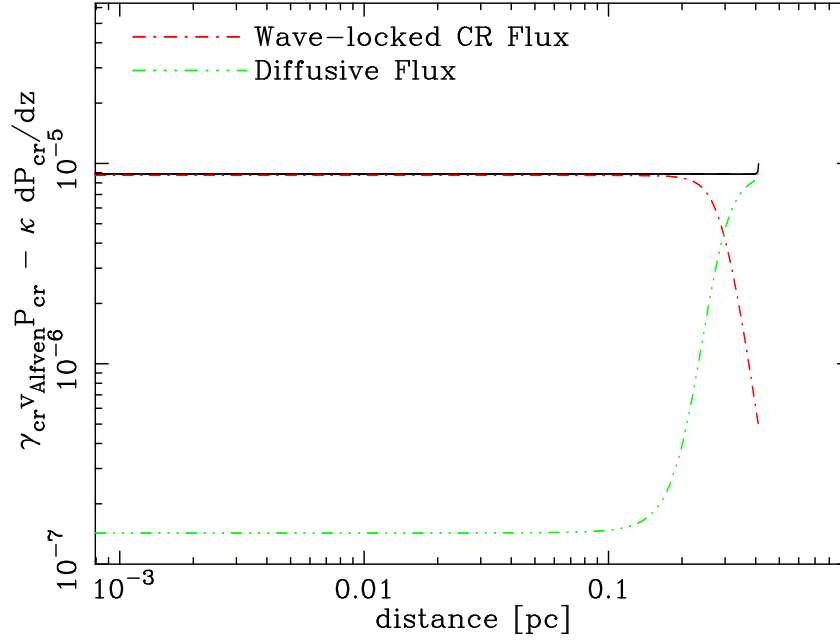


Figure 3. The transition from an advectively-dominated cosmic-ray flux to the diffusively-dominated flux. In summary, as the cosmic rays enter the cloud, the decrease in the Alfvén speed leads to a sharp drop in the advective flux which must be compensated (in the steady state) by an increase in the diffusive flux. Given the strong decrease in diffusivity, the cosmic-ray pressure is forced to drop in the cloud to yield the dominant diffusive flux. The slight up-tick in the diffusive flux at the very end of the above traces (at 0.4 pc) signals the face rise in diffusivity as ion-neutral damping becomes dominant, Alfvén waves die off, and the cosmic rays start free-streaming through the cloud.

With the above general picture outlined, we return to the equation governing cosmic-ray transport to explain the result more fully. It is easiest to look at a slightly expanded and simplified version of the cosmic-ray transport equation (Equation 2). If we write out that equation with the substitutions

$$\begin{aligned} \frac{d\kappa_{\text{cr}}}{dz} &= -\frac{\kappa_{\text{cr}}}{P_w} \frac{dP_w}{dz} \\ \frac{dv_A}{dz} &= -\frac{v_A}{2\rho} \frac{d\rho}{dz}, \end{aligned} \quad (20)$$

we get:

$$\gamma_{\text{cr}} \left(v_A \frac{dP_{\text{cr}}}{dz} + \frac{dv_A}{dz} P_{\text{cr}} \right) + \left(-\frac{d\kappa_{\text{cr}}}{dz} \frac{dP_{\text{cr}}}{dz} - \kappa_{\text{cr}} \frac{d^2 P_{\text{cr}}}{dz^2} \right) = (\gamma_{\text{cr}} - 1) \left(v_A \frac{dP_{\text{cr}}}{dz} + Q \right). \quad (21)$$

The first term in parentheses yields the change in the advective flux, while the second term in parenthesis gives the change in the diffusive flux. The terms on the right side of the equation give the various energy-loss mechanisms: the cosmic-ray pressure decreases due to wave generation, and also because of Coulomb interactions, pion productions, or ionization (these latter three processes symbolized by Q).

This equation helps one to see why the cosmic-ray pressure changes within clouds. First, we consider the term that describes the changes in advective flux. Since $dP_{\text{cr}}/dz < 0$, and $dv_A/dz < 0$ in the cloud (the Alfvén speed decreases as the density increases), both of these terms describe a *drop in the advective flux of cosmic-ray pressure into the cloud*.

Next, we consider the term that describes the change in diffusive flux. The first term in the second set of parentheses shows that the diffusive flux drops as the diffusivity drops (while dP_{cr}/dz is negative). Since $d\kappa_{\text{cr}}/dz = -\frac{\kappa_{\text{cr}}}{P_w} \frac{dP_w}{dz}$, as the wave pressure increases into the cloud (because of the increase in density; see Eq. 14), the diffusivity will decrease, and like the advective flux term, this component of the diffusive flux also drops. The second term then describes how the diffusive flux increases as $d^2 P_{\text{cr}}/dz^2$ becomes more negative.

As the cosmic-ray flux progresses into the cloud, the only self-consistent solution for the cosmic-ray pressure, in the steady-state, is for $d^2 P_{\text{cr}}/dz^2$ to become increasingly negative so that the cosmic-ray flux can become progressively diffusive and continue to propagate into the cloud. This explains why the cosmic-ray pressure starts to decrease into the cloud. This continues until the ion-neutral damping destroys the waves, or P_{cr} falls to zero (in extreme cases). *Beyond this point, cosmic rays are no longer locked to Alfvén waves when moving through the cloud, so no momentum or energy is communicated to the cool gas in the cloud.* Further, beyond this point, the cosmic-ray density in the cloud will be constant (neglecting cosmic-ray losses due to pion production, Coulomb interactions, or ionization). This result is consistent with previous results concerning cosmic-ray density

within clouds (e.g., Skilling 1971), although it conflicts with recent studies (Padoan & Scalo 2005), and addresses the cosmic-ray pressure in the cloud, which has not previously been explored.

4.3. Impact of Waves on the Cloud

The relatively strong cosmic-ray pressure gradient and the relatively high wave pressure on the outskirts of the cloud lead to heating on the cloud boundary as Alfvén waves are damped. Given the magnetic-field configuration that we have assumed (a constant magnetic field that directly connects the ambient medium to the cloud), we can compare the heating from wave damping to that from thermal conduction. Our calculations show that thermal conduction is dominant, as shown in Figure 4. To estimate the conductive heating and cooling, we use the thermal-conduction formulae of Pittard (2007): these estimates, for the fiducial model that we use here, show that the maximum conductive heating of the cool-cloud interface overpowers the maximum wave-damping by six orders of magnitude, although there are small regions where wave damping becomes comparable. This shows that cosmic-ray induced wave heating on the outskirts of the clouds is largely insignificant compared to the effects of thermal conduction. This is true even if we invoke a cooler ambient medium: if we modify the parameters of the ionized gas to keep the same pressure, setting $T = 10^4$ K and $n = 3 \text{ cm}^{-3}$, the Alfvén speed drops, so that the overall cosmic-ray flux is lower. In this regime, there is no rise in the cosmic-ray wave pressure on the outskirts of the cloud because of the relatively smaller rise in density, compared to the fiducial model. Therefore, there is less wave heating and the conductive heating still dominates in the cloud, even for a cooler ambient medium.

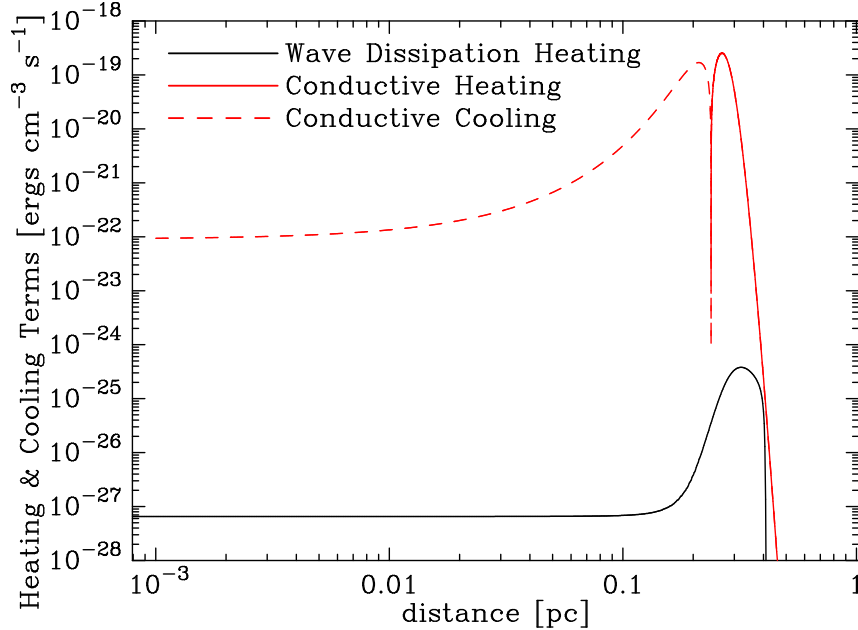


Figure 4. Comparison of the cosmic-ray derived wave-dissipation heating (given in the black solid line) to the impact of thermal conduction between the cool cloud and the ambient hot plasma, shown in red. Thermal conduction heating dominates by six orders of magnitude in the cool-cloud interface, and thermal conduction also more than compensates for the heating due to cosmic rays in the ionized plasma surrounding the cloud, as shown with the dashed-red line.

In regions where wave heating can grow to be competitive with conductive heating, do the Alfvén waves themselves become dynamically important? And do any of our test cases yield $P_w/P_B \gtrsim 1$, such that the wave growth becomes non-linear? At our highest magnetic-field strength, and at our highest cosmic-ray density, and in our cases of maximum wave heating, the maximum wave pressure is approximately 10% of the ambient thermal pressure; in our fiducial case, the wave pressure is only $\sim 10^{-4}$ of the thermal pressure. This shows that the wave pressure is not dynamically important, but may become important at larger cosmic-ray fluxes. Next, turning to the question of possible non-linearity in the wave growth, the maximum P_w/P_B that we find in our models is ~ 0.1 (and for our fiducial model, $P_w/P_B \sim 3 \times 10^{-4}$ at its peak); this shows that wave growth is still in the linear regime, although nearing the non-linear regime where $P_w/P_B \sim 1$. So, we do not see significant wave pressure or wave growth into the non-linear regime, although both of these concerns merit careful monitoring as cosmic-ray fluxes increase beyond those we consider.

4.4. Sensitivity to Initial Conditions & Assumptions

Throughout this work, we have assumed a particular form for the hot plasma/cloud interface, and a variety of parameters that describe both the wind and cloud structure. We have also specified boundary values for dP_{cr}/dz and P_{cr} . We check for the limits of applicability of our fiducial model, which showed that: (1) the cosmic-ray density does not change significantly within the cloud, (2) the cosmic rays are not coupled to gas within the cloud, and (3) cosmic-ray induced heating is insignificant compared to thermal conduction. We investigate changes in these results as we modify the cloud interface width, the cloud density, ionization

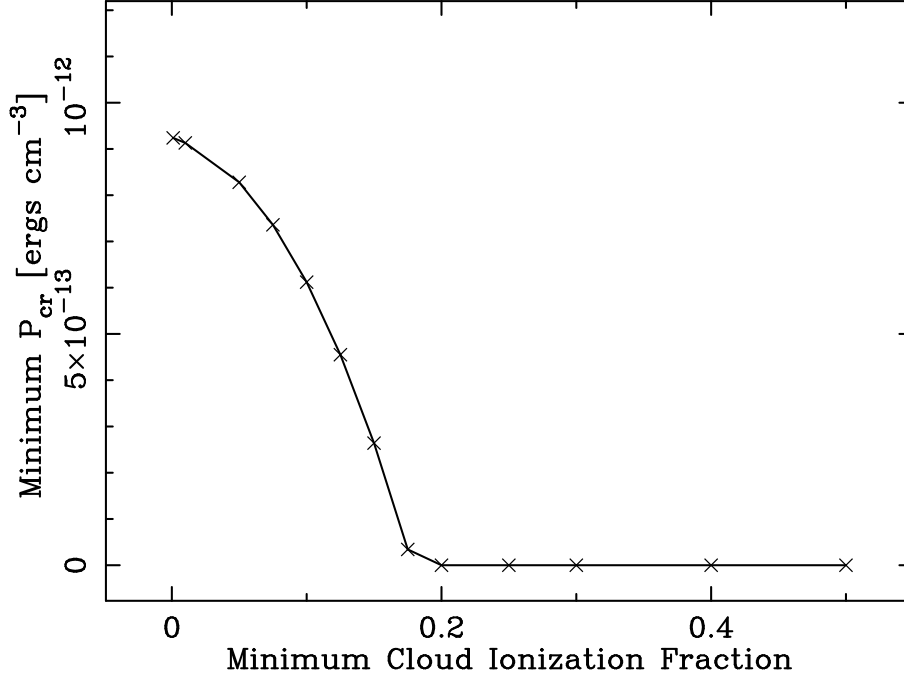


Figure 5. The minimum cosmic-ray pressure force as a function of the lowest level of ionization in the cloud.

profile, magnetic field strength, initial cosmic-ray pressure, initial cosmic-ray pressure gradient, cosmic-ray energy per nucleon (previously set to 1 GeV) and total cosmic-ray energy density. The parameters that we vary, and the range over which those parameters are varied, are summarized in Table 1.

4.4.1. Changes to Cosmic-Ray Pressure in the Cloud?

First, do changes to any of these parameters modify the fiducial-model result that cosmic rays do not couple to gas in the clouds? No. In all of our calculations, we see this transition in the boundary of the cool cloud. In all of our models, the cosmic-ray streaming instability fails at the cloud boundary (due either to ion-neutral damping or a severe drop in cosmic-ray pressure), and cosmic rays do not transfer energy and momentum to the gas within the clouds.

We do find changes in the cosmic-ray pressure on the cloud boundary, before the transition from the hydrodynamic limit to the kinetic limit (we refer to the last cosmic-ray hydrodynamic pressure within the wave-locked zone as the “minimum cosmic-ray pressure force”). None of the models show an increase in the cosmic-ray pressure in the cloud. The cosmic-ray pressure decreases slightly on the cloud boundary in the fiducial case, so one might expect that as we increase the penetration of wave-locked cosmic rays into the cloud, a larger drop in cosmic-ray pressure on the cloud boundary would be seen. This is true, although we always see a quick transition to the kinetic regime.

We find that the parameters that significantly affect the cosmic-ray pressure on the boundary of the cloud are the minimum cloud ionization fraction, the magnetic-field strength, and the initial cosmic-ray pressure. In the case of the minimum cloud ionization fraction, as the cool cloud’s minimum ionization level increases (from an ionization fraction of 10^{-3} to about 0.1-0.2), the onset of ion-neutral damping is delayed, and the cosmic-ray pressure decreases further, as Alfvén waves penetrate further into the cloud, before ion-neutral damping takes over. Also, for increased magnetic field strengths ($\sim 5 - 10 \mu\text{G}$), the advective cosmic-ray flux is higher than in the fiducial case (because of a higher Alfvén speed), and so the cosmic-ray generated waves dominate farther into the cloud than in the fiducial case before ion-neutral damping destroys the waves altogether. For $B > 10 \mu\text{G}$, dP_{cr}/dz is so strongly negative that P_{cr} drops to zero. The changes in the minimum cosmic-ray pressure are shown in Figures 5 and 6. In a similar way, increasing the initial cosmic-ray pressure also gives a stronger cosmic-ray flux at the surface of the cloud, which penetrates further into the cloud before ion-neutral damping takes over, and yields a change in the cosmic-ray pressure at the cloud interface (see Fig. 7).

4.4.2. Changes to Boundary-Layer Heating in the Cloud?

As shown in Section 4.3, conductive heating overall strongly dominates the heating from cosmic-ray generated Alfvén waves, although in very limited regimes, wave damping can compete with conductive heating (see Fig. 4). As we vary the parameters in Table 1, the heating at the cloud boundary changes, but in most cases and in most places on the cloud boundary, those changes yield heating rates that are orders of magnitude lower than that of thermal conduction. However, we have found that, as we increase the magnetic-field strength and the initial cosmic-ray pressure, there are small regions of the cloud boundary that can become dominated by wave heating. We present these results in Figures 8 and 9, which show the factors by which conductive heating dominates the wave heating when we change the magnetic-field strength and the initial cosmic-ray pressure, respectively. Given the varying strengths with distance of both the wave heating and conductive heating (see Fig. 4), comparing these rates

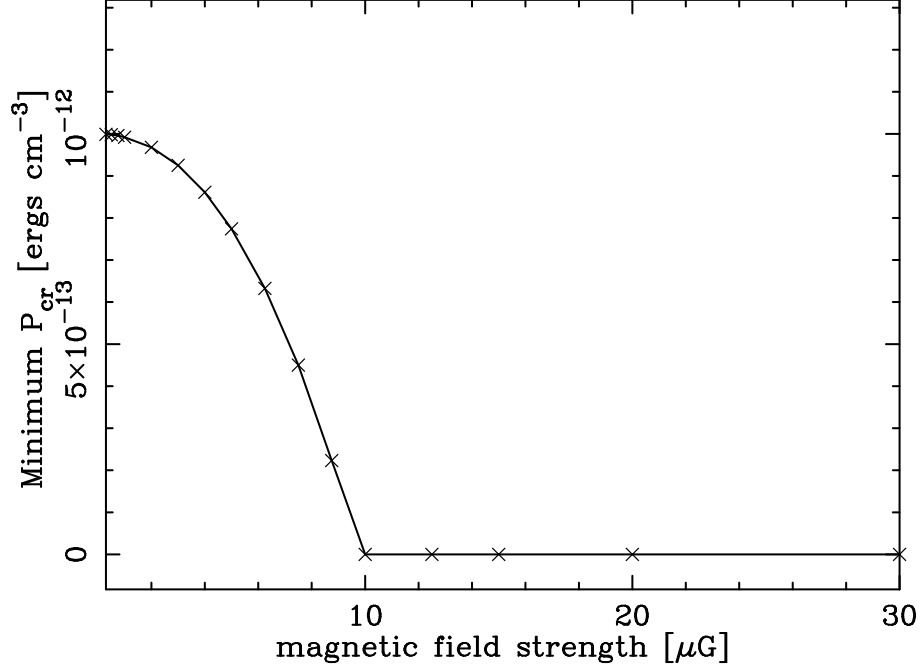


Figure 6. The change in the minimum cosmic-ray pressure force with an increase in the (uniform) magnetic field.

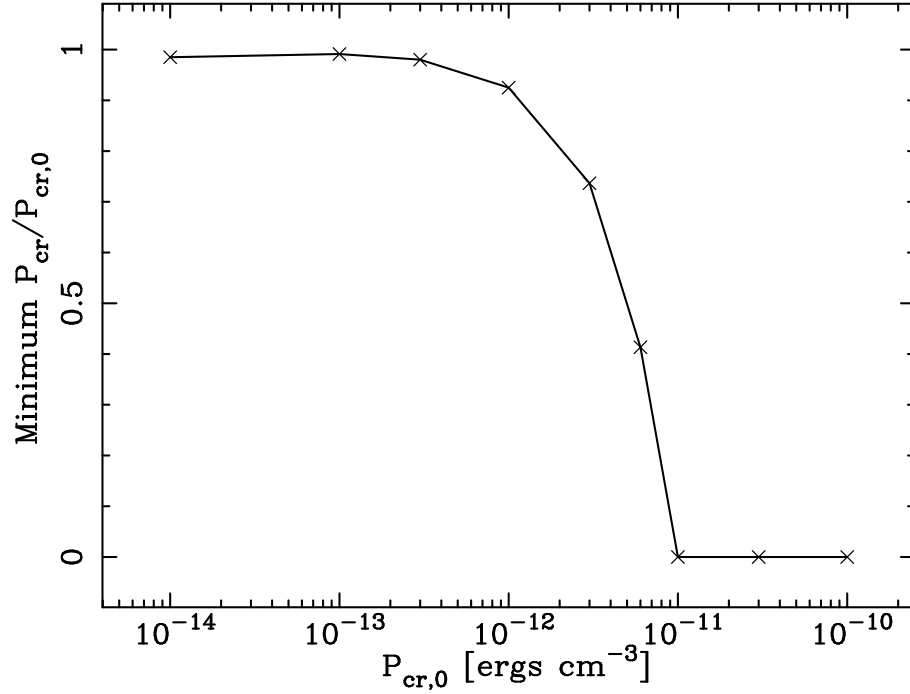


Figure 7. The minimum cosmic-ray pressure force in the cloud as a function of the initial cosmic-ray pressure in the wind ($P_{\text{cr},0}$). As the cosmic-ray pressure decreases, the effects of cosmic-ray pressure are felt to smaller and smaller depths into the cloud, with smaller and smaller differences between the external and internal cosmic-ray pressures force on the cloud boundary.

with a single ratio is too simplified, so we compute the ratio of conductive to wave heating as a function of distance, and plot both the maximum ratio of conductive heating to wave heating (solid line), and the minimum ratio (dashed line).

As one might expect from our discussion of cosmic-ray pressure changes in the previous section, when we increase the magnetic-field strength or increase the external cosmic-ray pressure, the heating from Alfvén-wave damping increases (see Fig. 8), at least over a limited region of parameters. For some conditions ($1.5 \times 10^{-12} \text{ ergs cm}^{-3} \lesssim P_{\text{cr},0} \lesssim 2.5 \times 10^{-11} \text{ ergs cm}^{-3}$ and $4 \mu\text{G} \lesssim B \lesssim 12 \mu\text{G}$), the magnetic-field strength and cosmic-ray pressures are high enough to yield a stronger cosmic-ray

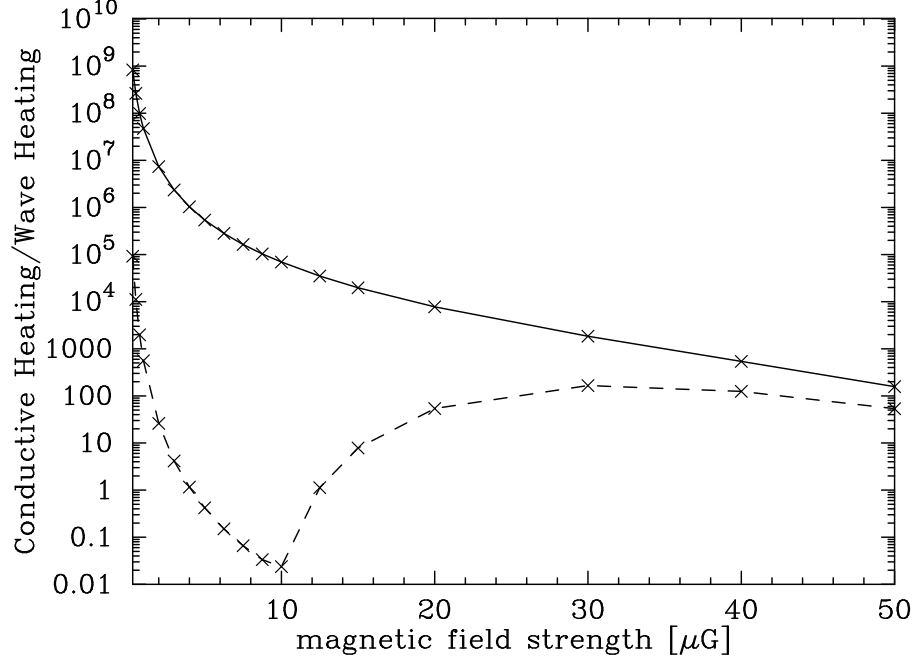


Figure 8. The maximum and minimum ratio of conductive heating to wave-damping heating, plotted as a function of increase in the (uniform) magnetic field strength; the minimum ratio is shown with a dashed line, the maximum ratio is shown with a solid line. As the magnetic-field strength increases, the Alfvén speed increases, and the flux of cosmic rays increases; this leads to a stronger wave pressure that extends deeper into the cloud, where, for intermediate field strengths of $4 \mu\text{G} \lesssim B \lesssim 12 \mu\text{G}$, wave damping can compete with conductive heating. Beyond that limit, the cosmic-ray pressure drops off too quickly into the cloud, and again, conductive heating dominates.

Table 1

The parameters and values we have assumed in the fiducial study and the variations in those parameters that have been examined to test robustness of our results.

Parameter	Fiducial Value	Range Tested
Cloud Interface Width (Δz_{edge})	0.05 pc	0.01 pc - 0.1 pc
Maximum Cloud Density (ρ_{max})	100 cm^{-3}	10 cm^{-3} - 1000 cm^{-3}
Cloud Minimum Ionization (x_{ion})	10^{-3}	0.01 - 0.5
Magnetic Field Strength (B)	$3 \mu\text{G}$	$0.3 \mu\text{G}$ - $30 \mu\text{G}$
Initial Cosmic-Ray Pressure ($P_{\text{cr},0}$)	$10^{-12} \text{ ergs cm}^{-3}$	$10^{-14} \text{ ergs cm}^{-3}$ - $10^{-10} \text{ ergs cm}^{-3}$
Initial Cosmic-ray Pressure Gradient ($\frac{dP_{\text{cr}}}{dz} _0$)	$-10^{-34} \text{ ergs cm}^{-4}$	$-10^{-31} \text{ ergs cm}^{-4}$ - $-10^{-39} \text{ ergs cm}^{-4}$
Cosmic Ray Energy per nucleon	1 GeV	1 MeV - 1 GeV

flux into the cloud (compared to our fiducial case), which can give a large Alfvén-wave amplitude in the region where conductive heating drops off, and in those regions, the wave heating can be greater than conductive heating. Towards the upper limit of our range of magnetic-field strength and cosmic-ray pressure, however, the cosmic-ray pressure drops so strongly on the boundary of the cloud (see Figs. 6 and 7) that the wave heating declines before it can begin to compete with the drop-off in conductive heating. Of course, these results may vary with other cloud configurations, but this analysis shows that it is possible for wave heating to compete with conductive heating in limited circumstances. It is also very important to note, though, that as shown in Figure 4, conductive heating is still dominant overall, despite the small regions in the cloud where wave damping might be competitive.

5. LIMITATIONS, COMPARISON TO PREVIOUS WORK, & IMPLICATIONS

This calculation makes a number of important assumptions, which we summarize here. First of all, we assume that there is a large-scale magnetic field which threads both the ionized gas and the cool cloud at constant strength. If the field strength changes, magnetic mirroring effects could also be important (e.g., Cesarsky & Völk 1978; Chandran 2000a), but we do not consider such effects here. This is a simplification, but one that is observationally warranted: Crutcher et al. (2010) find that the magnetic-field strength does not scale with density for $n \lesssim 300 \text{ cm}^{-3}$ (below the densities that are important in this work), and Heiles & Troland (2005) find that the magnetic-field strength in the cold, neutral medium of the ISM is of the same order as in other ISM components (this may perhaps be understood by ambipolar diffusion, see Zweibel 2002).

Also, we have considered only the role of the streaming instability in scattering cosmic rays; however, this is the dominant process for $\sim 1 \text{ GeV}$ cosmic rays (Kulsrud 2005). It is important to ask if other forms of turbulence may help scatter cosmic rays; however, turbulent wave modes, especially on the very small scales of the cosmic-ray gyroradius ($r_g \sim 10^{12} \text{ cm}$) are believed to lose energy very quickly in molecular clouds (e.g., Mac Low et al. 1998; Ostriker et al. 2001). In addition, those modes do not

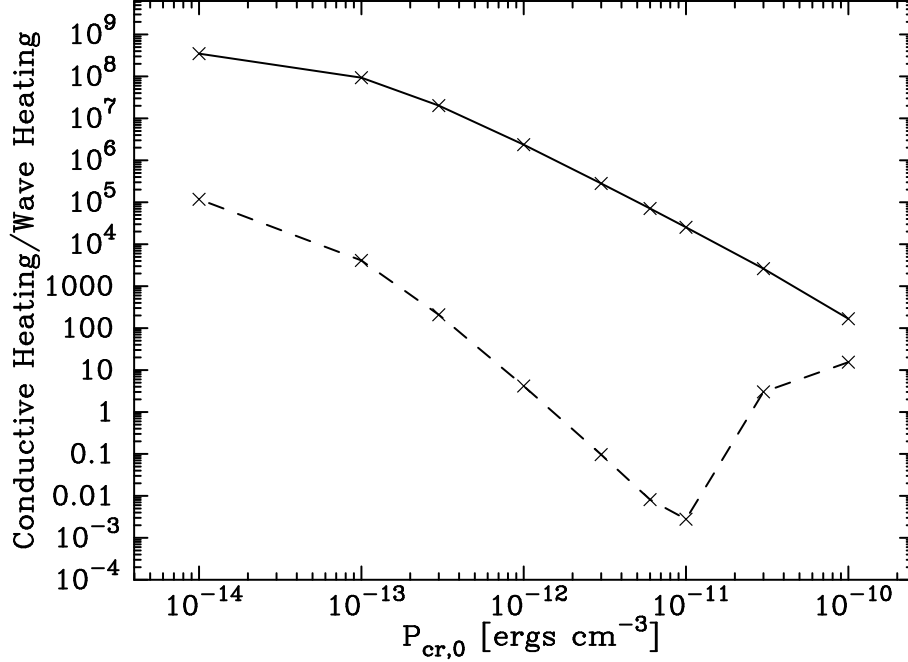


Figure 9. As in Figure 8, but as a function of increase in initial cosmic-ray pressure. As with the variation of magnetic-field strength, there is a drop in the minimum ratio at intermediate cosmic-ray pressures ($1.5 \times 10^{-12} \text{ ergs cm}^{-3} \lesssim P_{cr,0} \lesssim 2.5 \times 10^{-11} \text{ ergs cm}^{-3}$) due to the rise of the cosmic-ray flux, and then a drop at higher cosmic-ray pressures due to the strong drop-off in cosmic-ray pressure into the cloud.

scatter cosmic rays efficiently (e.g., Chandran 2000b; Yan & Lazarian 2004), so we do not expect the diffusivity to be affected by those processes.

Finally, when considering the role of cosmic rays in galactic winds, we do not consider the drop in velocity of the cool cloud as cosmic rays enter the cloud. Such a drop in velocity is likely to occur, as the hot-wind component is usually inferred to move at thousands of km/s (e.g., Strickland & Heckman 2009), and the clouds have velocities of $\lesssim 300\text{--}500$ km/s (e.g., Martin 2005; Steidel et al. 2010), so that the velocity difference could be of order a factor of four. In our calculations, though, the Alfvén speed already drops by an order of magnitude; a smaller drop in the bulk gas velocity would have the same, but smaller, effect on the transition of cosmic rays into the cool cloud. We infer that such a change in the bulk velocity will therefore only further increase the cosmic-ray flux injected into the clouds, and will have a similar impact on cosmic-ray pressure and heating as increases in the magnetic field strength and cosmic-ray pressure have. Therefore, the effects of increased velocity in the hot, ambient medium would not qualitatively change the results presented here.

How do our results compare to previous work? In the two sections below, we consider previous work on cosmic rays in cool clouds in the interstellar medium and in the galactic winds.

5.1. Cosmic Rays in the Multiphase ISM

The small drop in cosmic-ray pressure that we predict might seem similar to the “exclusion” of cosmic rays from clouds that was found by both Skilling & Strong (1976) and Cesarsky & Völk (1978). However, our model differs from those previous works in significant ways, and therefore this new model brings new effects to light. In those earlier papers, the exclusion of cosmic rays was driven by a drop in cosmic-ray density at $E \lesssim 300$ MeV due to ionization losses; as low-energy cosmic rays traversed the entire cloud, they were destroyed in collisions. It is also important that those works assumed a uniform cosmic-ray pressure outside of the cloud, and so assumed that the streaming instability only operated on the outskirts of the clouds.

In comparison, our work assumes an initial cosmic-ray pressure gradient (see Section 4.1), such that the streaming instability operates throughout the hot phase of the ISM. This assumption seems reasonable as we have found that even very small cosmic-ray pressure gradients lead to the growth of the streaming instability, which seems plausible since the streaming instability can explain the observed near isotropy of cosmic rays at the Earth (see, e.g., Kulsrud 2005). With this configuration, our models show a qualitatively new effect at higher energies: the transition of cosmic rays from the advective regime to the diffusively-dominated regime leads to a (small) cosmic-ray pressure drop in the outskirts of clouds. This effect is also present at low energies (down to $E = 1$ MeV), and is more important in our models than ionization losses (at least for cosmic rays just entering the outskirts of the clouds).

In the end, it is difficult to directly compare these models, but we believe that the assumption of zero cosmic-ray pressure gradient in the ISM (the assumption made in the papers of Skilling & Strong and Cesarsky & Völk) is not correct, and that therefore the “exclusion” of cosmic rays from cool clouds is not as strong as calculated in those papers (as we find in this work); in our models, we find only very small drops in cosmic-ray pressure within clouds. However, this cannot yet be a clear-cut comparison: we stress that the impact of low-energy cosmic-ray losses throughout the entire cloud (instead of just on the boundary, as we calculate here) is indeed important for large cloud columns as the previous papers have pointed out, and so

should be considered further. However, this realization points to a prediction: the observation of a *strong* decrease in the density of cosmic rays in cool clouds at low energy (by an order of magnitude or two at $E \sim 10$ MeV) would argue for the picture of Skilling & Strong (1976, see their Fig. 1) and Cesarsky & Völk (1978), whereas a relatively constant density of low-energy cosmic rays would argue for the model in this paper. Recent results that the cosmic-ray density does not change much or may even be higher than expected (e.g., McCall et al. 2003; Indriolo et al. 2010a) seem to point towards the model presented here, but more detailed studies are needed, both in the theory and perhaps in further observations. Overall, such studies would be useful as they may also allow insight into the presence or absence of a cosmic-ray gradient in the ISM, and therefore help us understand the propagation of cosmic rays in the ISM. (Of course, there are extremes of parameter space where the current model also yields drops in the cosmic-ray pressure within clouds, but these are rather extreme values; see Section 4.4.) We stress that both papers agree, however, that the density of cosmic rays is relatively constant in clouds at $E \sim 1$ GeV, which is the main area of focus in this work, and that, of course, cosmic-ray ionization losses will occur, and should be modeled in a complete picture of the full cloud for the most accurate comparison.

Putting aside those concerns about low-energy cosmic rays, we conclude that for most of the parameter space of the ISM, the cosmic-ray pressure does not change significantly before the cosmic rays become free-streaming. One might ask if the decreased transit time of cosmic rays that are not locked to Alfvén waves (with transit times of order $\sqrt{3}L_{\text{cloud}}/c$ vs L_{cloud}/v_A) might help decrease the gamma-ray emission, but since cosmic rays can be scattered back and forth across the cloud by Alfvén waves on either side of the cloud, it is not clear that the total residence time of cosmic rays inside clouds is significantly different from those for cosmic rays that are locked to Alfvén waves in a similar volume. Overall, at least, these results suggest that for clouds illuminated by cosmic rays and perhaps observed in gamma rays (e.g., Gabici et al. 2007; Torres et al. 2008; Abdo et al. 2010), the cosmic-ray density inferred within the clouds should not be significantly different than the density inferred outside of the clouds. We only predict decreases for the cosmic-ray density in clouds when the ambient medium has a very high cosmic-ray pressure (an order of magnitude higher than in the solar neighborhood) or when the (assumed uniform) magnetic-field strength is somewhat high (with values approximately twice the local field strength), or is more highly ionized than our fiducial case (minimum ionization fractions greater than ~ 0.1 lead to significant drops in the cosmic-ray density).

In contrast to these results, Padoan & Scalo (2005) inferred that cosmic-ray density would increase significantly within the cool clouds; such an increase in cosmic-ray density was predicted for clouds of density $\lesssim 100 \text{ cm}^{-3}$. We do not see this effect. It appears that the equations in Padoan & Scalo (2005) assume that the cosmic rays always travel at the local Alfvén speed in cool clouds. In contrast, the analysis presented in this paper shows that diffusivity quickly becomes dominant, and does not allow such a strong increase in cosmic-ray density.

5.2. Cosmic Rays in Multiphase Galactic Winds

As mentioned in the introduction, acceleration of cool clouds in winds without those clouds being quickly destroyed (on timescales of order a few times the sound crossing time in the external medium) is difficult to understand in clouds embedded in galactic winds. If cosmic-ray pressure had a smooth gradient within the cloud, perhaps it could input momentum to the cloud in a distributed and less disruptive way; however, we find that cosmic rays do not couple to gas within the body of the cloud. There is therefore no “body force” on the whole cloud due to cosmic rays; like ram pressure, cosmic rays would exert a pressure difference on the boundary that could perhaps help accelerate the cloud, but would add to the pressure differences already acting to disrupt the cloud. This may indicate another reason why radiative acceleration of clouds might be important (see, e.g., Murray et al. 2010), although perhaps any mechanism that imparts momentum to clouds will introduce instabilities (in the case of radiation pressure, see Mathews 1986). Connecting to our other main results, perhaps the heating on the outskirts of clouds by wave damping could help pressurize the surface of the cloud and slow down other instabilities. Meanwhile, an understanding of the mechanisms by which cosmic-ray density can drop in clouds will be important for predictions of gamma-ray emission.

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